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A New Mechanism of Kähler Moduli Stabilization in Type IIB Theory

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Abstract

We study the scalar potential in supersymmetric (orientifolded) Calabi Yau flux compactifications of Type IIB theory. We present a new mechanism to stabilize all closed string moduli at leading order in $\alpha^{'}$ by consistently introducing fluxes. As usual we consider the dilaton and the complex structure moduli stabilized by turning on three-form fluxes that couple to the F-part of the scalar potential. The Kahler moduli get fixed by the combined action of the flux-induced scalar masses with the magnetic fields of the open string sector, and Fayet-Illiopoulos terms. For supersymmetric three-form fluxes the model is $N = 1$, otherwise the mass terms are the scalar soft breaking terms of the MSSM fields. For the case of imaginary self dual three-form fluxes (ISD), the mass terms are positive and the minimum of the potential is at exactly zero energy. We argue that, under generic assumptions, this is a general mechanism for the full stabilization of closed string moduli. The vacua depend explicitly on the fluxes introduced in the manifold. A concrete realization of this mechanism for type IIB on a $(\frac{T^6}{Z_2 \times Z_2})$ orientifold is provided.

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1 Introduction

The problem of stabilization of moduli fields in string theory (scalar fields with flat directions in the potential) has been extensively studied, see for example [1, 2, 3, 4, 5, 6, 7, 8, 10, 11, 12, 13, 14, 15], since it has important theoretical and phenomenological consequences. From the experimental point of view the existence of such massless fields should be observed in fifth force experiments [1], but since it is not, the consistency of the theory requires such fields to acquire a mass. Cosmological observations also impose constraints on the moduli fixing scale to reproduce reheating at the inflationary epoch, having as a lower bound 100 TeV. The moduli fields come from two different sectors: closed and open string sectors. Moduli associated to the closed string sector give information about the size (Kähler moduli) and the shape (complex structure moduli) of the compact manifold, and also the dilaton. Moduli associated to the open string sector correspond to the presence of Wilson lines and to the parametrization of the position of D-branes in those cases when they are present, in the transverse dimensions.

A relevant result found in [3] showed that a linear combination of RR, and NSNS three-form fluxes on IIB theory were able to stabilize the dilaton and the complex structure moduli. This mechanism presented a great advance in the resolution of the problem although the Kähler moduli remained unfixed. The basic reason for this is due to the superpotential. It does not contain any dependence on the Kähler moduli, leading to a no-scale scalar potential at leading order in $\alpha^{'}$. In [4] KKLT found a way to fix one overall Kähler modulus by using non perturbative mechanisms such as condensation of gauginos or instanton effects. This led to an Anti De-Sitter (AdS) vacuum in a particular compactification manifold. By explicitly breaking supersymmetry through the introduction of anti-D3 branes and by fine tuning fluxes, they were able to lift it to a de Sitter vacuum. de Sitter vacua have acquired great importance due to the recent data that suggest the acceleration of the universe and also because of their close relationship with the inflationary scenario [5].

New advances in the context of the landscape have been achieved by Douglas et al. [6] obtaining manifolds with all moduli fixed through non-perturbative mechanisms and able to lead to a static cosmology. These last advances, although significant, have not been able to provide for realistic compactification manifolds. A potential problem in generating non-perturbative superpotentials from strong infrared dynamics is that it is model dependent. It can also generate too much massless charged matter.

[1] I would like to thank J. Conlon for this remark.

In [16] the model of KKLT was improved. They induce a supersymmetric model where one Kähler modulus is present by introducing magnetic fluxes contained on $D7$ branes.

Model building in this context has several fine-tuning and stability issues. In particular, the superpotential induced by the fluxes must be hierarchically small ($< 10^{-4}$) in order to obtain solutions with large volume in which the effective field theory approximation can be trusted. On the other hand the fluxes must fix the dilaton at small string coupling to suppress loop effects. Very recently some papers have appeared [7], see also [17], that find a way to stabilize all moduli by considering the combination of α'^3 effects and non perturbative contributions to the superpotential, giving rise to a large volume AdS vacuum. Its minimum is independent of the value of the flux superpotential. In this model supersymmetry is broken by the Kähler moduli and the gravitino mass is not flux dependent through the superpotential.

Here we propose a perturbative mechanism of moduli fixing which is fullfilled at the supersymmetric minimum of the theory and it is able to lead to realistic descriptions with all of the moduli fixed. In principle one could think that it is possible to induce any other dependence in the superpotential by introducing branes in the model. However it has been conjectured by [18, 19] that B-type branes (D-branes which wrap 2n-cycles with magnetic fluxes and the type of branes interesting in IIB models) can couple to the Kähler moduli only through Fayet-Illiopoulos (FI) terms. Several other works have also studied this problem. We will consider a IIB theory compactified on Calabi-Yau (CY) orientifolds, with RR, NSNS 3-form fluxes and magnetic fluxes. Three-form fluxes couple to the three cycles via the superpotential associated to the F-term, which does not depend on the Kähler moduli. We show that taking into account the coupling of the 3-form fluxes to the open string sector of magnetized D-branes which leads to flux induced mass terms, i.e.μ-terms or soft breaking terms with magnetic fields, together with FI terms gives a scalar potential which stabilizes the Kähler moduli. This mechanism is model independent and we think that it is a generic procedure for stabilizing moduli in a manifold. One advantage of this method is that by being supersymmetric or breaking spontaneously supersymmetry, it can be put in the 4D standard supergravity form and it keeps control over the types of interactions that can be induced. In a previous paper [20, 21] a method was proposed to stabilize some Kähler moduli through the coupling of fluxes to the FI-term in order to fix the blow-up moduli of the model. In that case the expected soft breaking term contribution did not include magnetic fields and failed in the attempt of stabilizing the full Kähler moduli. This idea is an extension of that one including magnetic fields in the configurations. In [22], in a different approach, they also consider mass terms as a mechanism to stabilize moduli.

In [23, 24] a type I theory was proposed with three-form fluxes and magnetic fluxes. They claim that they are able to stabilize Kähler moduli just through FI-terms. We argue that what they find does not constitute a true stabilization of the moduli since the moduli are free to acquire any other vev without any energy cost. To stabilize Kähler moduli it is necessary to also have a coupling of fluxes to the open string sector.

The paper is organized as follows: in section 2 we give a brief summary of three-form flux stabilization. We show how a linear combination of three form fluxes stabilizes complex structure moduli and the dilaton under very general assumptions. Non perturbative mechanisms are used to stabilize Kähler moduli. In section 3 we review how soft terms appear. We particularly focus on the flux induced soft breaking terms with magnetic fields and we find a general expression for these terms in a toroidal orientifold generalizing to the case when all Kähler moduli are different. In section 4 we describe D-term supersymmetry breaking. FI terms that couple to B-type branes represent a deviation from the supersymmetry conditions for the branes and a shift in the value at which the moduli have a supersymmetric value. However a proper recombination mechanism can restore the supersymmetry. In section 5 we propose the mechanism for stabilizing moduli without introducing non perturbative mechanisms and we describe in detail the minimization of the scalar potential. In section 6 we provide a concrete realization of this mechanism (of phenomenological interest) for the case of ISD three-form fluxes. We perform IIB compactified on $T^6/Z_2 \times Z_2 \times \Omega R$ moduli stabilization and we indicate the explicit values at which Kähler moduli get fixed. We conclude in section 7 with a discussion, summarizing the main results.

2 Review of moduli stabilization

In type IIB theory on a Calabi-Yau manifold the closed string moduli content associated to the geometry are: an axion-dilaton S, h_{11} Kḧaler moduli T_i parametrizing the CY_3 size and h_{21} complex structure moduli U_i parametrizing the CY_3 shape, where h_{11}, h_{21} are the Hodge numbers characterizing the Calabi-Yau three-fold. The Kähler potential associated to the closed string sector has, for toroidal compactifications where the metric factors are into three two-by-two blocks, the following expression,

$$k^2\mathcal{K}(S,U_i,T_i) = -ln[S+\overline{S}] - \sum_{j=1}^{h_{11}} ln[T_j+\overline{T}_j] - \sum_{j=1}^{h_{21}} ln[U_j+\overline{U}_j]. \tag{2.1}$$

Giddings, Kachru and Polchinski [3] introduced the methods to stabilize the dilaton and complex structure moduli by turning on 3-form fluxes. In type IIB theory, strings can have

RR and NSNS 3-form field strengths, which can wrap dual 3-cycles labeled by A and B leading to quantized background fluxes,

$$\frac{1}{4\pi^2\alpha'}\int_A F_3 = M \quad \frac{1}{4\pi^2\alpha'}\int_B H_3 = -K \tag{2.2}$$

where K and M are arbitrary integers. These forms allow consistent string compactifications of generic orientifold CY_3 manifolds. Fluxes also have some other interesting consequences: they induce a warp factor in the metric that deforms the manifold and generates hierarchies [20], fix the moduli partially, and also give a mechanism to break supersymmetry in a controlled manner by inducing soft breaking terms.

To understand the mechanism of the partial fixing of moduli we have to remark that these fluxes generate a superpotential found by [25]

$$W = \int_{CY_3} G_3 \wedge \Omega_3, \tag{2.3}$$

where $G_3 = F_3 - SH_3$, with S the complex axion-dilaton of type IIB theory. Ω_3 is the unique (3,0) form of the Calabi-Yau threefold. The holomorphic three form Ω_3 has a non-trivial dependence on the complex structure moduli U_i. This superpotential is independent of the Kähler moduli T_i. It gives the $D = 4, N = 1$ scalar potential

$$V_F = \exp^{k^2\mathcal{K}}(K^{I\overline{J}} D_I W \overline{D_J W} - 3k^2|W|^2) \tag{2.4}$$

Here I, J label all the above geometric moduli of the manifold.

The covariant derivatives are defined as $D_I W = \partial_I W + k^2 \partial_I \mathcal{K} W$ where $\mathcal{K}$ denotes the Kähler potential and $K^{I\overline{J}}$ the inverse of the Kähler metric defined in terms of the Kähler potential $K_{I\overline{J}} = \partial_I \partial_{\overline{J}} \mathcal{K}$.

This potential is of no-scale type ($\Lambda = 0$ at tree level in α') since the Kähler dependence on T_i cancels exactly the $3W^2$ contribution. Since the potential is positive definite, the global minimum of this potential lies at zero. The values of complex structure moduli and the dilaton get fixed for the particular values at which the superpotential is minimized, $D_i W = 0$, where i runs over all fields except the Kähler moduli. The Kähler moduli however remain unfixed as the superpotential has no dependence on them. The minimum of the potential is supersymmetric if $D_T W = \partial_I \mathcal{K} W = 0$ and non-susy otherwise.

The Kähler moduli have been stabilized by a different mechanism, namely non-perturbative contributions that introduce an exponential dependence. This contribution together with the one induced by fluxes gives the total superpotential

$$W = W_{flux} + A\exp(-aT) \tag{2.5}$$

and generate a scalar potential which typically has an AdS minimum at a finite value of T and a run-away behaviour at infinity. Several mechanisms for breaking supersymmetry allow in principle this minimum to be lifted to a de Sitter vacuum, i.e. [4, 16, 26].

3 Soft breaking terms with magnetic fields

The low energy effective action can have susy-breaking soft terms coming from magnetized or non magnetized configurations. Soft terms are operators of $dim < 4$ which do not induce quadratic divergences that spoil the good properties of the $N = 1$ supersymmetry. They give mass to the superpartners of the SM fields, when spontaneous supersymmetry breaking occurs. Their scale is expected to be not much above the electroweak scale. In fact when they are induced through fluxes they are of the order of the flux scale. These soft terms arise from the interaction of low dimensional D-brane charges induced on the D7 branes with the background flux (see [27, 28]). The soft terms contain gaugino masses, trilinear terms, and scalar masses of MSSM fields. Regarding the stabilization of Kähler moduli we are only going to be interested in scalar masses since they are the dominant terms in the expression. In this section we want to extend the results found in [29], (see also [30, 31]) to more general settings. Unless specified otherwise we will follow the notation of [29] but extending their results to be valid for a general configuration of D-branes compactified on toroidal orientifolds of type IIB theory on $T^6/Z_2 \times Z_2 \times \Omega R$.

We consider a six torus factorized as $T^6 = \otimes_{i=1}^{3} T_i^2$, performing the quotient by $Z_2 \times Z_2 \times \Omega R$ symmetry as in [11], filled with D7 branes some of which are magnetized and D9 branes carrying anti-D3 brane charges in the hidden sector to cancel RR and NSNS tadpole conditions while preserving $N = 1$ supersymmetry. The magnetic constant field is defined in terms of the wrapping number m of each stack of $D7_a$ branes transverse to the torus i, and n which represents the units of magnetic flux.

$$\frac{m_a^i}{2\pi} \int_{T_i^2} F_a^i = n_a^i \tag{3.1}$$

F_a^i is the world-volume magnetic field. For later convenience we define the following angles,

$$\Psi_a^i = arctan(2\pi\alpha' F_a^i) = arctan(\frac{\alpha' n_a^i}{m_a^i A_i}) \tag{3.2}$$

A_i represents the area of the two torus. Open strings give rise to charged fields. There are two types of states living on stacks of Dp-branes. `Untwisted` states are chiral fields coming

from open strings whose ends are attached to the same stack of branes and `twisted` sector chiral fields those lying at two different stacks of magnetized Dp-branes. We are only going to consider the twisted sector fields as they are the ones that have bosonic soft breaking terms. From D7 brane twisted sectors $D7_i - D7_j$ there are only chiral massless multiplets denoted by C^{7i7_j} transforming in the bifundamentals of $G_i \times G_j$, with G_i being the gauge group associated to the enhanced symmetry of each stack of D7 branes. The low energy dynamics of the massless fields is governed by a $D = 4, N = 1$ supergravity action that depends on the Kähler potential, the gauge kinetic functions and the superpotential. The moduli in this type of constructions, as already explained in [29], are $\{M, C_I\}$ with M the closed string moduli and $C_I = \{C^{7_i}, C^{7_i7_j}\}$, where $C_i^{7_i}$ are the moduli fields representing the position of the $D7_i$ branes in the transverse T_i^2 whereas $C_j^{7_i}$ correspond to the presence of Wilson lines turned on, in the two complex directions parallel to the $D7_i$ brane. In the following we will not care about open string moduli since there are models of phenomenological interest free of them, such as the one we propose in the last section. Closed geometric moduli that appear in the 4D $N = 1$ supergravity action consist of the complex dilaton

$$S = e^{-\phi_{10}} + ia_0 \tag{3.3}$$

where a_0 is the R-R 0-form and $e^{\phi_{10}} = g_s$ the string coupling constant; the complex structure moduli $U_j, j = 1, 2, 3$, which are equivalent to the following geometrical moduli for the case of toroidal compactifications,

$$\tau_j = \frac{1}{e_{jx}^2}(A_j + ie_{jx}.e_{jy}) \tag{3.4}$$

where e_{jx}, e_{jy} are the T_j^2 lattice vectors, A_j is the area of the two-torus, which for the particular case of the square T_j^2 is equal to $A_j = R_{jx}R_{jy}$, and the dual angle is $\varphi_a^i = \arctan(\frac{n_a^i R_{ix}}{m_a^i R_{iy}})$. The geometric Kähler moduli ρ_i are described by,

$$\rho_j = A_j + ia_j, \tag{3.5}$$

where the axions a_j arise from the R-R 4-form. However, as explained in [29, 30, 31] the Kähler moduli field denoted by $T_i, i = 1, 2, 3$ are not equivalent to the geometric moduli ρ_i, since its correct expression is

$$T_i = \exp^{-\phi_{10}} \frac{A_j A_k}{\alpha'^2} + ia_i \quad j \neq i \neq k \tag{3.6}$$

and the gravitational coupling in $D = 4$ is $G_N = k^2/8\pi$ with

$$k^{-2} = \frac{M_{pl}^2}{8\pi} = \exp^{-2\phi_{10}} \frac{A_1 A_2 A_3}{\pi\alpha'^4}. \tag{3.7}$$

The string scale is defined as $M_s = \frac{1}{\sqrt{\alpha'}}$. We have chosen this type of compactification for its simplicity: The moduli in this type of compactification are just the dilaton S, three complex structure moduli associated to the relation between the radius of each tori $U_i, i = 1, ., 3$ and three Kähler moduli associated to the size of each torus $T_i, i = 1, 2, 3$.
The Kähler potential contains the Kähler part $\widehat{K}$ associated to the closed string moduli, $M = \{S, T_i, U_i\}$, and the Kähler part $\widetilde{K}$ which is associated to the matter fields $C_I = \Phi_{aa}, \Phi_{ab}$ of the open string sector [29].

$$k^2\mathcal{K} \;=\; k^2\widehat{K} + k^2 \sum_{IJ} \widetilde{K}_{IJ} C_{\overline{I}} C_J + k^2 \sum_{IJ} Z_{IJ}(C_I C_J + c.c) + \dots \tag{3.8}$$

The standard expression for the soft breaking terms found in [32] is

$$m_I^2 = m_{3/2}^2 + V_0 - \sum_{M,N} \overline{F}^{\overline{M}} F^N \partial_{\overline{M}} \partial_N log(\widetilde{K_{IJ}}) \tag{3.9}$$

in supergravity conventions with all quantities measured in Planck units. $I = aa, ab, ba, bb$ labels the different stacks of magnetized D7 branes. F^M are the auxiliary fields of the corresponding chiral multiplet Φ_A which in general have the following expression

$$\overline{F}^{\overline{A}} = k^2 \exp^{k^2 K/2} K^{\overline{A}B} D_B W \tag{3.10}$$

where $\Phi_A = \{M, C_I\}$ and as before $C_I = \{C_j^{7_i}, C^{7_i 7_j}\}$. When a spontaneous breaking of supersymmetry occurs, the auxiliary vevs acquire a vacuum expectation value. In this case the supersymmetry breaking is produced by the presence of a non supersymmetric 3-form flux, (G_3 containing a $(0, 3)$ piece). The auxiliary fields are parametrized as

$$F^S = \sqrt{3} C s m_{3/2} \sin(\theta) \exp^{-i\gamma_s}, \tag{3.11}$$

$$F^{T_i} = \sqrt{3} C t_i \eta_i m_{3/2} \cos(\theta) \exp^{-i\gamma_i}, \tag{3.12}$$

such that $\sum_i \eta_i^2 = 1$ and η and θ, the goldstino angle, control whether S or T_i dominate the SUSY breaking. Here

$$C^2 \;=\; 1 + \frac{V_0}{3m_{3/2}^2} \tag{3.13}$$

$$m_{3/2} \;=\; \exp^{1/2k^2\mathcal{K}} W \tag{3.14}$$

$$V_0 \;=\; F^m \widehat{K}_{mn} F^n - 3m_{3/2}^2 \tag{3.15}$$

and $m_{3/2}$ denotes the gravitino mass and V_0 the cosmological constant.

For the case of soft breaking terms induced by 3-form and magnetic fluxes, the expression becomes

$$\begin{aligned} m_I^2 \; &= m_{3/2}^2 + V_0 - 1/4\overline{F}^{\overline{M}} F^N \partial_{\overline{M}} \partial_N \ln(st_1t_2t_3) + \textstyle\sum_{i=1}^3 (\partial_{\overline{M}} \partial_N \nu_i) \overline{F}^{\overline{M}} F^N (ln(u_i)) \\ &\quad -1/2 \textstyle\sum_{i=1}^3 \overline{F}^{\overline{M}} F^N \partial_{\overline{M}} \partial_N (\ln(\Gamma(1-\nu_i)) - \ln(\Gamma(\nu_i))). \end{aligned} \tag{3.16}$$

where the Kähler potential has the following form,

$$k^2\mathcal{K} \;=\; k^2\widehat{K} + k^2 \sum_{IJ} \widetilde{K}_{IJ} C_{\overline{I}} C_J + k^2 \sum_{IJ} Z_{IJ}(C_I C_J + c.c) + \ldots \tag{3.17}$$

$$\widetilde{K} \;=\; \frac{(st_1t_2t_3)^{\frac{1}{4}}}{2\pi\alpha'} \Pi_{j=1}^3 u_j^{-\nu_j} \sqrt{\frac{\Gamma(1-\nu_j)}{\Gamma(\nu_j)}} \tag{3.18}$$

$$k^2\widehat{K} \;=\; -\ln(s) - \sum_i \ln t_i - \sum_i \ln(u_i) \tag{3.19}$$

with the conventions of [29]

$$s = S + \overline{S}; \quad t_i = T_i + \overline{T_i}; \quad u_i = U_i + \overline{U_i}. \tag{3.20}$$

$k^2 = 4\pi\alpha'(st_1t_2t_3)^{-1/4}$, and $\widehat{\nu}_i = \frac{1}{\pi}(\Psi^i_{ab}) = \frac{1}{\pi}(\Psi^i_b - \Psi^i_a)$, where $\Psi^i_a = \arctan(g_a\beta_i)$ with $\beta_i = \sqrt{\frac{st_i}{t_jt_k}}$ and $g^i_a = \frac{n^i_a}{m^i_a}$. $\sum \widehat{\nu}_i = 0$ trivially, which is the condition for two stack of D-branes to preserve the same supersymmetry. The ν_i are computed in terms of $\widehat{\nu}_i/\pi$, such that $0 < \nu_i < 1$ and $\sum_{i=1} \nu_i = 2$ as in [29]. Here $\nu_i = 1 + \widehat{\nu}_i/\pi$ iff $\widehat{\nu}_i \leq 0$ and $\nu_i = \widehat{\nu}_i/\pi$ otherwise. This last expression is closely related to the one in [29], although it has been generalized to the case when all of the Kähler moduli are different.

Regrouping terms,

$$\begin{aligned} m_{ab}^2 \; &= m_{3/2}^2 + V_0 - 1/4\overline{F}^{\overline{M}} F^N \partial_{\overline{M}} \partial_N \ln(st_1t_2t_3) + \textstyle\sum_{i=1}^3 (\partial_{\overline{M}} \partial_N \nu_i) \overline{F}^{\overline{M}} F^N \\ &\quad (ln(u_i) - 1/2 B_0^i(\nu_i)) - 1/2 \textstyle\sum_{i=1}^3 B_1^i(\nu_i) \overline{F}^{\overline{M}} F^N (\partial_{\overline{M}} \nu_i \partial_N \nu_i). \end{aligned} \tag{3.21}$$

The B_p are defined in terms of polygamma functions as,

$$B_0^i(\nu_i) \;=\; \frac{\partial_\nu \Gamma(1-\nu_i)}{\Gamma(1-\nu_i)} - \frac{\partial_\nu \Gamma(\nu_i)}{\Gamma(\nu_i)}, \tag{3.22}$$

$$B_p^i \;=\; \partial_\nu B_{p-1}(\nu_i), \tag{3.23}$$

and in terms of a useful analytical expression is

$$B_0(z) = \pi cotan(\pi z) \tag{3.24}$$

and its derivatives. Let us remark that this definition of $B_p(\nu)^i$ is slightly different of the one used in [29]. A special value of this function is at $\nu_i = 1/2$ where $B^i_k(1/2) = 0$.

We finally obtain the following expression,

$$\begin{aligned} m^2_{ab} &= m^2_{3/2} + V_0 - 3/4C^2 m^2_{3/2}[1 + 1/4\pi^2 \sum_{i=1}^{3}(ln(u_i) - 1/2B^i_0(\nu_i)) \qquad (3.25) \\ & \quad [4\mathbf{P}_i \sin(2\pi\Psi_i)_{ab} - \mathbf{Q}_i \sin(4\pi\Psi_i)_{ab}] + 1/8\pi^3 \sum_{i}^{3} B^i_1(\nu_i)\mathbf{Q}_i(\sin(2\pi\Psi_i)_{ab})^2] \end{aligned}$$

where $\sin(\pi\Psi_i)_{ab} = \sin(\pi\Psi_i)_b - \sin(\pi\Psi_i)_a$. and the above variables $\mathbf{P}, \mathbf{Q}$ are defined in terms of goldstino angles,

$$\mathbf{P}_i = \sin^2\theta + \cos^2\theta(\eta_i^2 - \eta_j^2 - \eta_k^2) \qquad (3.26)$$

$$\begin{aligned} \mathbf{Q}_i &= 1 - 2\cos^2\theta\{\eta_i\eta_j cos(\gamma_i - \gamma_j) + \eta_i\eta_k cos(\gamma_i - \gamma_k) - \eta_j\eta_k cos(\gamma_j - \gamma_k)\} \qquad (3.27) \\ & \quad + \sin(2\theta)\{\eta_i cos(\gamma_i - \gamma_s) - \eta_j cos(\gamma_j - \gamma_s) - \eta_k cos(\gamma_k - \gamma_s\} \end{aligned}$$

so they are sensitive to the particular choice of 3-form fluxes.

The expression for soft breaking mass terms can be finally expressed as,

$$m^2_{ab} = \frac{1}{4}(1 + 3A)m^2_{3/2} + \frac{1}{4}(3 + A)V_0 \qquad (3.28)$$

where

$$A \equiv -1/4\pi^2 \sum_{i=1}^{3}(ln(u_i) - 1/2B_0(\nu_i))[4\mathbf{P}_i \sin(2\pi\Psi_i)_{ab} - \mathbf{Q}_i \sin(4\pi\Psi_i)_{ab}] \qquad (3.29)$$

$$-1/8\pi^3 \sum_{i}^{3} B_1(\nu_i)\mathbf{Q}_i(\sin(2\pi\Psi_i)_{ab})^2. \qquad (3.30)$$

These results recover the ones of [29] by making appropriate substitutions for their particular configuration and imposing $t_2 = t_3$. Although the Kähler part of the potential $\widetilde{K}_{C\overline{C}}$ correspond to the twisted magnetized sectors, with the above definitions, the soft terms at twisted, unmagnetized $D7_2 - D7_3$ stacks are also correctly found due to appropriate cancellations.

4 D-term supersymmetry breaking

In this section we are interested in characterizing D-term behaviour in the presence of B-type branes. The motivation is the following: in order to solve the problem of Kähler moduli

stabilization, naively we can think of finding a perturbative generalization of the superpotential that contains Kähler moduli in its definition. However this seems not to be possible. Kähler moduli are the moduli associated to the $(1,1)$ forms that naturally can be thought of as being stabilized by magnetic fluxes living on D-branes. In general there are h_{11} of them, being the h_{11} Hodge number counting the number of 2-cycles present in the compactified manifold. The Decoupling Statement of Douglas e al. [18, 19] establishes that Kähler moduli on Type IIB can only couple to B-type branes (i.e. Branes wrapping even cycles) through Fayet Iliopoulos (FI) terms in the scalar potential. The mirror statement of this on type IIA says that complex structure moduli can only couple to A-branes (branes wrapping odd cycles (3-cycles)) through FI terms. This fact seems disappointing from the perspective of finding a perturbative superpotential for Kähler moduli. Moreover this is a common feature for the case of D-branes at singularities [18, 19].

In [35, 36] an $\mathcal{N}=1$ 4D type IIA theory on $\frac{T^6}{Z_2\times Z_2}$ orientifold with D6 branes at angles was studied. In this section we will review the main properties. This model is dual to magnetized D9 branes on type IIB theory with discrete torsion. Requiring supersymmetric models imposes a condition between orientifold planes and D branes. On type IIA, in a supersymmetric model each stack of D6 branes is related to the orientifold planes O6 by a rotation in $SU(3)$. The supersymmetry configurations imposes a condition on the angles θ_i of the D6 branes with respect to the orientifold plane in the i-th direction of the two-torus, which for the case they considered was,

$$\sum_{i=1}^{3}\theta_i = 0. \tag{4.1}$$

This condition is equivalent to

$$\sum_{i=1}^{3}\arctan\left(\chi_i\frac{m_i}{n_i}\right) = 0, \tag{4.2}$$

with R_{xi}, R_{yi} the radius of the i-th two-torus and $\chi_i = \frac{R_{yi}}{R_{xi}}$ the untwisted complex structure moduli. This model is T-dual to a type IIB orientifolded theory with discrete torsion and twisted Kähler moduli with the condition

$$\sum_i \Psi_a^i \;=\; \frac{3\pi}{2}mod(2\pi), \tag{4.3}$$

where we are using the convention for angles introduced in the preceding section, $\Psi_a^i = arctan(\frac{\alpha' n_a^i}{m_a^i A_i})$ in terms of the tori areas A_i. A_i are expected to couple open string modes

in the D9 branes on 2-cycles (B-branes) as Fayet-Illiopoulos (FI) term. This is the dual version to the one indicated in [35]. In [39] it is explained that FI terms are proportional to the deviation from the susy condition, giving an effective action proportional to the deviation,

$$\sum_a \int dx^4 \xi_a D_a \tag{4.4}$$

where for small FI terms , [40], $\xi_a = \delta \sum_i \Psi_a^i = (\sum_i \Psi_a^i - susy \quad cond.)$ and vanishes for supersymmetric configurations of D-branes. Since $D_a = \sum_b^{N_b} |\phi_{ab}|^2 - |\phi_{ba}|^2 + \xi_a$, and $V_D = \frac{1}{2} D_a D^a$ then, the D-term piece of the scalar potential is then

$$V_D = \sum_a^{N_a} (\sum_b^{N_b} |\phi_{ab}|^2 - \sum_b |\phi_{ba}|^2 + \xi_a)^2, \tag{4.5}$$

where $|\phi_{ab}|$ represents the charged matter fields lying on an oriented string with ends attached at two different stacks a and b. $|\phi_{ba}|$ represents the matter fields at the intersection of the same two stacks a and b with reversal orientation and has also to be considered. N_a, N_b are the number of parallel Dp-branes at each stack.

The dependence of the scalar potential on a ϕ_{ab} matter field is

$$V_D(\phi_{ab}) = (|\phi_{ab}|^2 + 1/2\xi_a)^2 + (-|\phi_{ab}|^2 + 1/2\xi_b)^2 \tag{4.6}$$

where we have renamed for the second part of the R.H.S $a \to b$ and we have used that $|\phi_{ba}|^2 = -|\phi_{ab}|^2$ and the mixed term which gives a mass term for $|\phi_{ab}|$

$$|\phi_{ab}|^2 (\xi_a - \xi_b) = |\phi_{ab}|^2 \delta\Psi_{ab}, \tag{4.7}$$

defining $\delta\Psi_{ab} = \xi_b - \xi_a$ and approaching $(\delta\Psi_{ab})^2 \sim \frac{1}{2}(\xi_a^2 + \xi_b^2)$, we arrive to the familiar expression for D-terms appearing from twisted sector, [40]:

$$V_D = \sum_I (q_I |C_I|^2 - \delta\Psi_I)^2 \tag{4.8}$$

where I runs over all the indices of the matter field , i.e, aa, ab, ba, bb. q_I are the charges of the matter fields under the $U(1)$ gauge group of the D7 branes, with $q_{aa} = 0, q_{ab} = -q_{ba}, q_{ab} = +1, 0, -1$; and $C_I = \phi_{aa}, \phi_{ab}$ are all of the open string moduli. The FI terms appear always when there are supersymmetry breaking coming from the 2-form magnetic fluxes. In standard type IIB orientifolded actions the FI terms can be properly tuned by adjusting the twisted moduli. Physically this term reproduces at leading order the splitting between scalar and fermion masses [35]. Namely in the ab sector, chiral fermions remain massless at tree level while their scalar partners obtain a mass proportional to $\delta\nu_{ab} = \frac{1}{\pi}\sum_i(\Psi_a^i - \Psi_b^i)$.

An important remark regarding supersymmetry is the following: In [35] it has been pointed out that since some of these scalar fields can acquire vevs, the existence of a FI terms by itself does not automatically imply a susy breaking since they may acquire them so as to make the D-term vanish. Physically it is due to a recombination process of the D6 branes -in type IIA picture- (which now are not supersymmetric since the angles have changed) into supersymmetric smooth 3-cycles. This process gives a vev to some of the scalar fields ϕ_{ab}, at the intersection.

In the type IIB picture, the argument remains valid. D9 branes have magnetic fields at supersymmetric values. In the presence of non-vanishing FI terms the D9 branes become non susy, but if some scalars acquire mass cancelling D-term contribution, then they change their magnetic values on the 2-cycles restoring supersymmetry.

It has also been argued that although the existence of a FI-term in local compactifications allows for supersymmetry breaking, for the case of global compactifications recombination processes seem to be a generic mechanism to restore supersymmetry.

As already mentioned in the introduction, in [23, 24] a mechanism was proposed to stabilize Kähler moduli through three-form fluxes and magnetic fluxes. The argument was that the Kähler moduli get fixed at its supersymmetric value since the supersymmetry condition for magnetized D-branes depends explicitly on them. However, because of the argument above, this is not a true stabilization since matter fields can acquire a vev without any energy cost to cancel the D-term contribution, modifying the value of the twisted moduli. Then the scalar potential has a flat direction with $V = 0$.

5 A new mechanism of moduli stabilization

We are going to consider a decoupling approach where complex structure moduli and the dilaton have been stabilized by the standard mechanism. Turning on suitable three-form fluxes stabilizes their vevs and allows their dynamics to be integrated out as in [4]. This approach is valid in the regime when the mass scale of the Kähler moduli is much less than that of the dilaton and the complex structure moduli. The warping induced by the fluxes will be negligible in the approximation of large volume. We propose the following mechanism to stabilize the Kähler moduli perturbatively: We consider the coupling of fluxes (3-form fluxes and magnetic ones) to the open string sector that induce bosonic flux induced masses. These mass terms in principle can be soft breaking terms or supersymmetric. Supersymmetric

mass terms are the μ terms appearing in the superpotential when supersymmetric 3-form fluxes (2,1) are turned on and the D-brane stacks share at least one parallel direction [31, 38]. These flux induced masses m_I^2 for magnetized D-branes combined with the FI contribution represent a new coupling in the scalar potential that lift the flat directions of the potential giving masses to all the moduli. This fact will happen generically when soft masses are present but as remarked above, for supersymmetric masses to exist on susy flux-backgrounds, only very specific configurations will be allowed. The effective scalar potential in $D = 4$ taking into account the F-piece and the D-piece of the potential is equal to

$$V_T = V_F^{back} + \text{flux induced mass terms} + FI = \tag{5.1}$$

$$= V_F^{back} + \sum_I m_I^2 |C_I|^2 + \sum_I (q_I |C_I|^2 - \sum_i \delta\Psi_I^i)^2, \tag{5.2}$$

where V_F^{back} is the background no-scale potential induced by the superpotential. Flux induced mass terms have the effect of lifting the flat directions in the potential. These terms can be positive or negative. For the case of ISD G_3 the mass terms are positive and V is also positive definite. We will see that there exists a global minimum of the potential V that fixes all of the moduli. In order to minimize this potential we should consider the minimization with respect to every moduli C_I, M. It is important to ask whether the resulting critical point is a saddle point or a minimum. The answer is that iff there exist a critical point such that $V_T = 0$, being V_T definite positive, then this point is a global minimum of the theory and in this case it corresponds to its supersymmetric value. (The supersymmetric minimum in fact corresponding to F-flatness and D-flatness condition) since it is bounded from below to a value greater or equal than 0. We find the value of this minimum at $|C_I| = 0, \sum_i \delta\Psi_I^i = 0$.

Let see it in more detail. To see the stabilization we need to minimize the potential with respect to the full moduli C_I, M, since the flat directions were associated in [35] to the presence of matter fields which did not acquire a vev. As explained before, the non-vanishing of the FI term represents a shifting from the supersymmetric condition associated to the D-brane configuration $\sum_i \delta\Psi_I^i \neq 0$. As we are interested in stabilizing Kähler moduli on type IIB, then we are going to impose that the flux induced terms have to be associated to the presence not only of 3-form fluxes but also of magnetic fields which depend explicitly on the Kähler moduli. The presence of magnetic fluxes gives an extra dependence of the intersecting angles $\sum_i \delta\Psi_I^i$ on the Kähler moduli.

The superpotential is not renormalized at any order in perturbation theory and receives no α' corrections from the bulk, however it could receive from brane contributions [37]. The

D-term, by an appropriate selection of fluxes, can leave the model supersymmetric and so will not receive corrections changing its structure. Moreover, if we analyse the equation we can see that this holds generically given $m_I^2 > 0$ (as we will see this is the case for the interesting ISD 3-form fluxes).

We have in principle two kinds of flux contributions, ISD fluxes and IASD fluxes. We do not consider interaction terms between them. Since IASD fluxes do not solve type IIB 10D equations of motion unless non-perturbative effects will be taken into account, we will not consider them in the analysis. In the presence of ISD fluxes only, as pointed out in [28], soft terms are positive. The argument relies in the following : they can be regarded as geometric moduli of the F/M-theory fourfold and generate positive definite scalar potential. In [41] it is argued that this is a general property of ISD fluxes also valid in the case of magnetic fluxes at the intersections [2]. A way of illustrating this point is that for the case of soft breaking terms with magnetic fluxes, the superpotential

$$W = \int_{CY_4} G_4 \wedge \Omega_4 \tag{5.3}$$

leads to a positive definite scalar potential. W in particular includes D7-brane geometric moduli [34], which in this case are described in terms of the homological charges also coming from the magnetized D-branes.

For the case of ISD fluxes, the $V_F^{TT} = 0$ since it is a no-scale potential ($D_T W \overline{D_T} W = 3\,|W|^2$), then extremizing with respect to the moduli T_N gives the equation

$$\partial_N V = 0 \rightarrow \sum_I (-2q_I |C_I|^2 \partial_N \Psi_I + 2(\delta\Psi_I)\partial_N \Psi_i + (\partial_N m_I^2)|C_I|^2) = 0. \tag{5.4}$$

Minimizing with respect matter field gives as solution

$$|C_I| \;=\; 0, \quad V(C_I) = |\delta\Psi_I|^2 \tag{5.5}$$

$$|C_I| \;=\; \pm\sqrt{\frac{2q_I\delta\Psi_I - m_I^2}{2q_I^2}}, \quad V_D(|C_I^2|_{min}) = \frac{m_I^2}{4q_I^2}(4q_I(\delta\Psi_I) - m_I^2). \tag{5.6}$$

The first solution imposes in this model that the global minimum lies at $\delta\Psi_I = 0$, fixing the moduli through the dependence of Ψ_I to its supersymmetric value iff soft terms > 0 for all I. The value of the minimum of the potential corresponds for supersymmetric configuration of D-branes, to a no-scale potential ($V = 0$), and in those cases when supersymmetry can be consistently broken through the FI term to a de Sitter minimum, in the same spirit as [16].

[2] We thank L. Ibanez for helpful clarifying explanation at this point

The other two extrema lead to an AdS vacua and they are only possible in those cases when $|C_I|$ is real, that correspond to have negative squared mass terms for the supersymmetric case, $m_I^2 < 0$, so it is not possible for ISD G_3 fluxes or to have $0 < m_I^2 < 2\sum_I q_I \delta\Psi_I$ for a non trivial FI term. Assuming that $\sum_I q_I \delta\Psi_I > 0$ one can see that $m_I^2 \geq 2\sum_I q_I \delta\Psi_I$ always. The bound is never saturated unless $\nu = 1/2, u_i = 1, s \sim 10^{-3}$ [3] and these values do not allow to have consistent flux compactifications on this background compatible with tadpole cancellation conditions, so these other two possibilities are never achieved nor by making fine tuning of fluxes. The unique minimum is at $C_I = 0$.

To be sure that critical points are not saddle points if we were interested in these cases associated to IASD fluxes, (that we are not), we should perform the Hessian calculation -very involved due to the higly non-trivial structure of the V_T- , including also the F-part of the potential that is no longer vanishing.

One could also ask whether trilinear couplings of the soft terms could change this behaviour. The answer is not for the particular case we are considering, i.e. ISD fluxes. let see it in more detail. The full structure of the soft and susy terms is rather complicated. Generically,

$$\text{soft terms} = \frac{1}{2}\sum_a (M_a \lambda^a \lambda^a + h.c.) + \sum_I m_I^2 |C_I|^2 + \tag{5.7}$$

$$\sum_{I,J,K} A_{IJK} C_I C_J C_K + 1/2 B_{IJ} C_I C_J + h.c. \tag{5.8}$$

where λ^a represents gaugino mass terms, m_I are the scalar mass terms and there are bilinear and trilinear couplings each of which with a highly nontrivial dependence on all of the moduli. However one can see that all of the induced mass that appear in the scalar potential are polynomial of lower bound 2 in the matter fields C_I, hence, as before $C_I = 0$ is still an extrema, and as before $\delta\sum_I \Psi_i = 0$ corresponds to $V_T = 0$ which is the global minimum of the theory. Moreover, the expression for the trilinear terms in this case corresponding to $W = W_{flux}$ is the following,

$$A_{IJK} = F^M \widehat{K_M} - \partial_M log(\widetilde{K_{II}} \widetilde{K_{JJ}} \widetilde{K_{KK}}) \tag{5.9}$$

trilinear terms can be estimated as $A_{IJK} \sim 1/2(m_{3/2} - \sum_i 1/\nu_i) \leq 1/2 m_{3/2}$ and scalar mass terms $m_I^2 \sim m_{3/2}^2(1/4 + 3/4(-1/\sum_i 1/\nu_i + \sum_i \frac{1}{\nu_i^2})) \geq 1/4 m_{3/2}^2$. As before one can ask if there exists an appropiate fine tuning in the configuration of D-branes and in the choice of 3-form fluxes in such a way that for a particular configuration of brane case the other two minima

[3] I would like to thank G. Tasinato for his comments regarding this point.

will be $C_I = 0$ with a $FI \neq 0$, giving $V_T = 0$ and the answer is that again it seems not possible. Both bounds are saturated for $\nu_i = 1/2$ which corresponds to the case cosidered above that leads to non physical solution in these set-ups. Scalar soft terms are the dominant contribution. An important clue in this result is the fact that the whole mass terms are positive defined for ISD G_3. Otherwise there could also other solutions corresponding to have $V_T < 0$, and in those cases a careful analysis should be performed.

In the absence of magnetic fluxes moduli can not be fixed since Ψ_I is independent of them. However as already shown in the preceding section the presence of magnetic fluxes does not by itself imply the stabilization of Kähler moduli, as we already explained, as matter fields are free to acquire any vev to cancel D-term contribution leaving Kähler moduli unfixed. We want to emphasize that only when there are magnetic fluxes in the 3-form flux induced soft (susy) [4] breaking terms combined with the FI term, the Kähler moduli get fixed. This contribution coupling to the FI-term, prevents matter fields from acquiring a vev cancelling this contribution as it is energetically disfavoured.

This is then a generic mechanism, in the same way that 3-form fluxes H_3, F_3 serve to stabilize the complex structure moduli $(2,1)$ and the dilaton through the superpotential. The magnetic fluxes given by a particular configurations of the magnetized D-branes fix the Kähler moduli (1,1) at their supersymmetric values, only once the potential has no flat directions. These directions are lifted by the combined action of these 3-form flux-induced soft terms with magnetic fields, and the FI-terms. We think that this solution gives a final answer to the question of perturbative stabilization of moduli. Both types of moduli need to be present to achieve a model without moduli. The scale of stabilization of the moduli vevs is at supersymmetric values. The mass scale of the moduli however depends on the model considered, and is given by the scale of the flux induced terms.

All of the previous discussion in principle remains valid for the case of susy flux induced terms, i.e. mass terms with the same masses for fermions and their scalar superpartners (it has been obtained in D-brane action calculations, see for example, [27], and F-theory [33, 34].). These susy mass terms, however do not appear always that there are three form fluxes ISD (2,1) [31], it is needed a D-brane configuration with at least two chiral multiplets such that they are at least parallel in one direction to generate μ-terms. To have a true $N = 1$ model it is also needed masses for all of the moduli present in the configuration, allowing to obtain at the same time non-trivial areas. If such a model is provided then it constitutes a model with

[4]This mechanism is also valid for supersymmetric masses generated by similar mechanisms whenever appropiate configurations are considered.

Kähler moduli stabilized. This implies that for a supersymmetric model, whenever the induced flux mass terms are all positive, which is the case for ISD fluxes, the matter fields are going to be fixed at zero. Spontaneous susy breaking will lift this value to the lowest metastable de Sitter vacua once the Kähler moduli have been stabilized by the supersymmetric condition. In both cases (non-susy and susy) this result is very appealing since it is valid for the ISD fluxes which are the ones we are interested in since they solve the 10D equations of motion in type IIB theory. From the calculational point of view, they stabilize the Kähler moduli for a given Dp-brane content without the need to specify a detailed expression for the highly complicated soft terms. We provide an explicit example in the last section of the paper.

5.1 Beyond imaginary self-duality condition

In general for IASD fluxes or a combination of both, the contribution of V_F has to be taken into account but this does not change the behaviour of the stabilization. The two possible extra extrema are

$$|C_I|^2 = \frac{2q_I\delta\Psi_I - m_I^2}{2q_I^2} \tag{5.10}$$

For a susy model this reduces to:

$$V(|C_I|^2_{min}) \quad = \quad -\frac{|m_I|^4}{4q_I^2} < 0 \quad iff \quad m_I^2 < 0 \quad \forall I. \tag{5.11}$$

The minima of the potential correspond to a AdS vacua. For the case of negative bosonic soft terms only very rough approximations to the stabilized values can be done because of the complicated equation to solve. Moreover the Kähler potential cannot be simplified so much since the bilinear Z_{ij} has to be calculated. This case can appear when IASD fluxes are taken into account. The explicit value at which Kähler moduli get fixed now it is going to be determined by the explicit expression of the flux induced mass terms, that one for the soft terms (see section 4) is highly nonlinear. On the other hand IASD fluxes do not solve the equations of motion in this description since it is purely done at perturbative level (the situation changes if non-perturbative mechanisms are taken into account), and they generate poorly understood run-away potentials. We will not perform the calculation.

6 An example of IIB on $\frac{T^6}{Z_2\times Z_2\times\Omega R}$

In this section we provide a concrete example of phenomenological interest which realizes the new mechanism proposed in the preceding section. Some examples of flux compactifications of

phenomenological interest are [45, 46]. See also [47], in which some of the flux compactifications of phenomenological interest are able to lead to KKLT models.

In the following we will review the main features of the constructions of [11, 43] which are a global embedding of the local model proposed in [44] and used in [29]. We will then construct our explicit realization.

The model of [11, 43] is based on type IIB string theory compactified on a $\frac{T^6}{Z_2 \times Z_2}$ modded out by the orientifold action, which has also been examined in[42]. We consider $T^6 = \Pi_{i=1}^3 T_i^2$. The generators of the orbifold symmetries $Z_2 \times Z_2$ are θ, ω which act on the complex coordinates of the tori as

$$\theta : (z_1, z_2, z_3) \to (-z_1, -z_2, z_3), \tag{6.1}$$

$$\omega : (z_1, z_2, z_3) \to (z_1, -z_2, -z_3). \tag{6.2}$$

The orientifold action is given by ΩR where Ω is the usual world-sheet parity and
$R : (z_1, z_2, z_3) \to (-z_1, -z_2, -z_3)$. This model contains 64 $O3$ planes each one on a fixed point of $\mathcal{R}$ and 4 $O7_i$ planes, located at the Z_2 fixed points of the i-th T^2 and wrapping the other tori. We consider the case of intrinsic torsion as in [46]. The open string sector contains D9 branes with non trivial magnetic fluxes. The non-trivial gauge bundle generically reduces the rank of the group and upon KK reduction leads to D=4 chiral fermions. The magnetic fluxes also induce D-brane charges of lower dimension that contribute to the tadpoles. Magnetized D9 branes usually have D7, D5 and D3 charges. As explained in [42] the topological information of these models is encoded in three numbers. $(N_a, (n_a^i, m_a^i))$ where N_a is the number of D9 branes contained on the a-stack, m_a^i is the number of times that a-stack of D-branes wrap the i-th T^2 and n_a^i is the units of magnetic flux in that torus induced by D-branes. The unit of magnetic flux of the D-branes in that torus, as seen in section 3, is

$$\frac{m_a^i}{2\pi} \int_{T_i^2} F_a^i = n_a^i. \tag{6.3}$$

The $D9_a$ branes preserve the same supersymmetry of the orientifold planes provided that

$$\sum_{i=1}^{3} \Psi_a^i = \frac{3\pi}{2} \quad mod(2\pi), \tag{6.4}$$

$$\text{with} \quad \pi\Psi_a^i = arctan(\frac{n_a^i \beta_i}{m_a^i}), \tag{6.5}$$

where no summation over index i is performed. This condition also guarantees that any two sets of branes preserve a common supersymmetry since the relative angles

$$\theta^i_{ab} = \Psi^i_b - \Psi^i_a \tag{6.6}$$

trivially satisfy

$$\sum_i \theta^i_{ab} = 0 \quad mod(2\pi). \tag{6.7}$$

In the case of a global analysis we need to add new branes which are expected to be in a hidden sector in order to cancel global tadpoles [42]. The conditions are,

$$1) \quad \sum_\alpha N_\alpha n^1_\alpha n^2_\alpha n^3_\alpha + 1/2 N_{fluxes} = 16 \quad N_{min} = 8 \quad \text{with torsion} \tag{6.8}$$

$$2) \quad \sum_\alpha N_\alpha m^i_\alpha n^j_\alpha m^k_\alpha = -16 \quad i \neq j \neq k \quad \text{and} \quad i, j, k = 1, 2, 3 \tag{6.9}$$

with $N_{flux} = 64.n \quad n \in Z$. Global cancellation of $\mathbf{Z_2}$ RR charges must also be imposed to cancel the contribution of $D5_i - \overline{D5_i}$ and $D9_i - \overline{D9_i}$ pairs, which is equivalent to satisfying the following conditions for a case with torsion,

$$\sum_\alpha N_\alpha m^1_\alpha m^2_\alpha m^3_\alpha \in \mathbf{4Z} \tag{6.10}$$

$$\sum_\alpha N_\alpha n^i_\alpha n^j_\alpha m^k_\alpha \in \mathbf{4Z} \quad i \neq j \neq k \quad \text{and} \quad i, j, k = 1, 2, 3 \tag{6.11}$$

and for the case without torsion the conditions are the same by making these changes,

$$\sum_\alpha N_\alpha n^1_\alpha n^2_\alpha n^3_\alpha + 1/2 N_{fluxes} = -16 \quad N_{min} = 4 \tag{6.12}$$

$$\sum_\alpha N_\alpha m^1_\alpha m^2_\alpha m^3_\alpha \in \mathbf{8Z} \tag{6.13}$$

$$\sum_\alpha N_\alpha n^i_\alpha n^j_\alpha m^k_\alpha \in \mathbf{8Z} \quad i \neq j \neq k \quad \text{and} \quad i, j, k = 1, 2, 3 \tag{6.14}$$

Let us focus on the case with torsion. The following model is a concrete realization that serves our purpose.

MSSM

$\mathbf{N}_\alpha$	(n^1_α, m^1_α)	(n^2_α, m^2_α)	(n^3_α, m^3_α)	$(\Psi^1_\alpha, \Psi^2_\alpha, \Psi^3_\alpha)$
$N_a = 8$	(1,0)	(g,1)	(g,-1)	$(\frac{\pi}{2}, \pi\delta_2, \pi - \pi\delta_3)$
$N_b = 2$	(0,1)	(1,0)	(0,-1)	$(0, \frac{\pi}{2}, \pi)$
$N_c = 2$	(0,1)	(0,-1)	(1,0)	$(0, \pi, \frac{\pi}{2})$
$N_{h_1} = 2$	(-4,1)	(-2,1)	(-3,1)	$(\pi(1-\varphi_1), \pi(1-\varphi_2), \pi(1-\varphi_3))$
$N_{h_2} = 2$	(0,1)	(-5,1)	(-4,1)	$(0, \pi(1-\phi_2), \pi(1-\phi_3))$
$8N_f$	(1,0)	(1,0)	(1,0)	$(\frac{\pi}{2}, \frac{\pi}{2}, \frac{\pi}{2})$

(6.15)

This includes a slight modification to the proposal of [46] since their configuration was not able to fix all of the Kähler moduli. In fact the spectrum keeps its interesting phenomenological properties except for the number of families which depend on the particular homological charges chosen. To satisfy the consistency conditions the following equation has to be solved,

$$4n + g^2 + N_f = 8. \tag{6.16}$$

It gives different flux vacua for the different possible replication of families $I_{ab} = \Pi_i(n^i_a m^i_b - n^i_b m^i_a)$. Searching for a realistic scenario it is necessary to obtain three family replication, unfortunately a simple examination reveals that this model does not contain it. The possible vacua are:

For $g = 2$

$$n = 0 \quad N_f = 4 \tag{6.17}$$

$$n = 1 \quad N_f = 0 \tag{6.18}$$

For $g = 1$

$$n = 0 \quad N_f = 7 \tag{6.19}$$

$$n = 1 \quad N_f = 3 \tag{6.20}$$

To illustrate our example we choose the vacua for $g = 1$, containing fluxes, i.e. $n = 1, N_f = 3$.

Spectrum

Locally the spectrum of this model can contain the MSSM since g is not constrained. This has been studied in [44]. This sector is called the visible sector. We show it below as a remainder,

Intersection	Matter fields	Rep	Q_{B-L}	Y
$D7_1 - D7_2$	Q_L	$3(3,2)$	1	1/6
$D7_1 - D7_3$	U_R	$3(\overline{3},1)$	-1	$-2/3$
$D7_1 - D7_3$	D_R	$3(\overline{3},1)$	-1	1/3
$D7_1 - D7_2$	E_L	$3(1,2)$	-1	1/2
$D7_1 - D7_3$	E_R	$3(1,1)$	1	-1
$D7_1 - D7_3$	N_R	$3(1,1)$	1	0
$D7_2 - D7_3$	H	$(1,2)$	0	1/2
$D7_2 - D7_3$	$\overline{H}$	$(1,2)$	0	$-1/2$

(6.21)

Table 1: Chiral spectrum of the MSSM-like model.

However the global completion imposes constraints in such a way that $g \neq 3$ so it is not possible to obtain 3-family replication and in that sense, we are proposing is a toy model although it keeps the nice properties of chiral matter, correct quantum numbers, etc.. It contains the intersection of the so-called, visible sector and the hidden sector and hidden-hidden sector. We analyze the full spectrum in terms of a Pati-Salam model.

Sector	Matter	SU(4)xSU(2)xSU(2)xUSp[24]	Q_a	Q_{h_1}	Q_{h_2}	Q'
(ab)	F_L	$(4,2,1)$	1	0	0	1/3
(ac)	F_R	$(\overline{4}$,1,2)	-1	0	0	-1/3
(bc)	H	(1,2,2)	0	0	0	0
(ah_1)		$6(\overline{4},1,1)$	-1	-1	0	5/3
(ah_1')		$4(\overline{4},1,1)$	-1	-1	0	5/3
(ah_2)		18(4,1,1)	1	0	-1	-5/3
(ah_2')		20(4,1,1)	1	0	-1	-5/3
(bh_1)		12(1,2,1)	0	-1	0	2
(bh_1')		12(1,2,1)	0	-1	0	2
(ch_1)		8(1,1,2)	0	-1	0	2
(ch_1')		8(1,1,2)	0	-1	0	2
(h_1h_1')		288(1,1,1)	0	-2	0	4
(h_1h_2)		12(1,1,1)	0	1	1	0
(h_1h_2')		196(1,1,1)	0	1	1	0
(fh_1)		(1,1,1)x[24]	0	-1	0	2
(fh_2)		(1,1,1)x[24]	0	0	-1	2

(6.22)

Table: Chiral spectrum of a one generation Pati-Salam $N = 1$ chiral model of table 1. The abelian generator of the unique massless U(1) is given by $Q^{'} = \frac{1}{3}Q_a - 2(Q_{h_1} - Q_{h_2})$.

Some linear combinations of U(1) will leave $U(1)$ fields massive in a Green-Schwarz mechanism but since our purpose is to show the perturbative stabilization of moduli we have not included them in the spectrum calculation. We show that a realization of string compactification on a CY_3 orientifold with some phenomenological properties render the moduli fixed.

When we substitute for the supersymmetry conditions,by taking the values of table (6.15) in (6.4), we obtain the following equations

$$\delta_2 = \delta_3 \to t_2 = t_3, \tag{6.23}$$

$$\pi(\varphi_1 + \varphi_2 + \varphi_3) = \frac{3\pi}{2}, \tag{6.24}$$

$$\pi(\phi_2 + \phi_3) = \frac{\pi}{2}. \tag{6.25}$$

Using (6.5) a straightforward calculation gives

$$\beta_1 = 0.157, \quad \beta_2 = \beta_3 = \frac{1}{\sqrt{20}} \tag{6.26}$$

which means that the vacuum expectation value at which Kähler moduli gets fixed is,

$$t_1 = 20s, \quad t_2 = t_3 = 127.4s, \tag{6.27}$$

where s denotes the vacuum expectation value of the dilaton which is the string coupling constant. Substituting in the value of areas (3.6) this means that

$A_1 = 6.38\alpha^{'}$	$A_2 = A_3 = 4.47\alpha^{'}$

(6.28)

This mechanism implies that for this toroidal compactification, are needed three and only three stacks of D-branes with different magnetic fluxes in such a way that lead to different equations in order to have a determined compatible system of equations. This is the subtlety that does not allow us to use Marchesano et al.'s model in our mechanism as an example of D-brane configuration since they have just two different equations. This construction is in no way unique, and represents a toy model to illustrate the mechanism of stabilization. We expect that more complicated models can be constructed with realistic spectrum and with all of the moduli fixed.

Regarding the blow-up moduli associated to the orbifold fixed points, they are going to be also stabilized with this mechanism since for models with intrinsic torsion, the moduli are Kähler, and can get fixed in the same way as it was done in [20]. With the above results we can perform the explicit calculation of soft mass terms for a given three-form flux configuration with complex-structure moduli and dilaton stabilized. We choose the one [12] for $n = 1, N_{flux} = 64$ with $\mathcal{N} = 0$ supersymmetry

$$G_3 = 2(d\overline{z}_1 dz_2 dz_3 + dz_1 d\overline{z}_2 dz_3 + dz_1 dz_2 d\overline{z}_3 + d\overline{z}_1 d\overline{z}_2 d\overline{z}_3), \tag{6.29}$$

$$W = 8(u_1 u_2 u_3 - s). \tag{6.30}$$

The complex-structure moduli and the dilaton are stabilized at $u_i = s = i$, which in our notation is equivalent to $u_i = s = 1$. This values lead to a non-acceptable value for perturbative analysis since $g_s = 1$ as explained in [43]. However we will use this to give an example of explicit calculation of the soft terms. The soft terms from T-dominance (ISD) fluxes correspond to $V_0 = 0$. The gravitino vev is

$$m_{3/2}^2 = \frac{|W|^2}{s\Pi_i t_i u_i} \tag{6.31}$$

and we have made $C = 1, cos(\theta) = 1, \eta_i = 1/\sqrt{3}, \gamma_i = \gamma_T$ as is explained in detailed in [29].

The soft breaking terms for this model are:

$m_{ab}^2 = 375, 2m_{3/2}^2$	$m_{ac}^2 = 7, 8m_{3/2}^2$	$m_{bc}^2 = \frac{1}{4}m_{3/2}^2$
$m_{ah_1}^2 = 120, 2m_{3/2}^2$	$m_{ah_2}^2 = 58, 5m_{3/2}^2$	$m_{af}^2 = 16, 9m_{3/2}^2$
$m_{bh_1}^2 = 79, 9m_{3/2}^2$	$m_{bh_2}^2 = 222, 3m_{3/2}^2$	$m_{bf}^2 = \frac{1}{4}m_{3/2}^2$
$m_{ch_1}^2 = 125, 6m_{3/2}^2$	$m_{ch_2}^2 = 225, 2m_{3/2}^2$	$m_{cf}^2 = \frac{1}{4}m_{3/2}^2$
$m_{h_1h_2}^2 = 244, 8m_{3/2}^2$	$m_{h_1f}^2 = 39, 4m_{3/2}^2$	$m_{h_2f}^2 = 30, 5m_{3/2}^2$

(6.32)

in terms of gravitino mass $m_{3/2}$. Since the microscopic source of SUSY-breaking is the above ISD flux in a toroidal setting, by using the above definitions, $|W|^2 = 256$, and the gravitino mass is $m_{3/2} = 0.019$, so the soft terms have unrealistic values as they are extremely high.

For general toroidal/orbifold models with intersecting D6-branes the flux-induced soft terms are typically of the order of the string scale $(1/\alpha')$, which is only slightly smaller than M_{Pl} and not able to solve hierarchy problems [29]. This fact is due to the simplicity of the compactification manifolds as well as the fact that fluxes are distributed uniformly.

7 Discussion and conclusions

We have shown a new method to dynamically stabilize, with fluxes, all moduli in supersymmetric and non supersymmetric models. In supersymmetric models we have explained that the mechanism is restricted to particular configurations able to generate μ terms for all of the moduli, i.e. configurations of stacks parallel to each other in at least one direction in which a suitable $(2,1)$ G_3 flux has been turned on. Three-form fluxes generically stabilize the dilaton and the complex structure moduli. They generate an F-term part of the scalar potential. Magnetic fluxes are two forms that fix Kähler moduli at their supersymmetric value once the potential has no flat directions. This goal is achieved by inducing through 3-form flux a non-susy (susy) breaking (flux-induced) mass terms with magnetic fields in the scalar potential, that combined with the FI term lift the flat directions. We think this is a generic mechanism that gives a final answer to the problem of perturbative moduli stabilization. These mass terms avoid the possibility of cancelling the D-term by the consistent adjustment of the matter fields vev, since it requires an energy cost. To leading order, the F-part of the potential is a non-scale potential and together with D-term are responsible for the stabilization of the Kähler moduli. ISD fluxes induce positive mass terms and fix the Kähler value of the D-term to its supersymmetric value. In the supersymmetric case it can be possible to implement a similar mechanism of the one found in [16] to generate a de Sitter space (by spontaneous supersymmetry breaking). However, possibly toroidal models will be unable to generate adequate fine tuning to guarantee that the different approximations are still valid (small cosmological constant, small quantum corrections, a big potential barrier). Maybe this mechanism combined with the one of [16] could be implemented in a more complicated model (it is necessary that $U(1)$'s of the FI term will not be charged or the matter will have some special properties that are not present in the case considered). For the case $m^2 < 0$ both parts of the potential, F- and D-, including the supergravity potential $V_F^{background}$ contribute and an explicit calculation in terms of the particular soft breaking terms is needed. They can be associated to IASD or a combination of (IASD and ISD) fluxes. The analysis of IASD contribution has not been performed as they do not lead to solutions of physical interest. The scalar potential generically has two AdS minima.

Clearly the case with $m^2 > 0$ which correspond to turning on ISD three form fluxes is much more interesting since it solves the equation of motion and stabilizes the moduli at a value of the order of the string scale which means that it does not depend on the scale of supersymmetry breaking and is higher enough to induce reheating processes at early stages of

the universes. The mass scale of the moduli presumably is of the order of the soft breaking mass scale (which is of the order of flux scale $\frac{\alpha'}{R^3}$). We have given a concrete Kähler moduli-free realization of phenomenological interest of Type IIB on $\frac{T^6}{Z_2\times Z_2\times\Omega R}$. However this model represents a toy model. We expect improvements in the search for realistic compactifications moduli-free (i.e. three family generation, lower soft masses) in more complicated scenarios as those with warped metrics due to throats, that are also able to explain the hierarchy problem, as well as, address inflation.

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